

Dark matter and a suppression mechanism for neutrino masses in the Higgs triplet model

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Abstract

We extend the Higgs triplet model so as to include dark matter candidates and a simple suppression mechanism for the vacuum expectation value (v_Δ) of the triplet scalar field. The smallness of neutrino masses can be naturally explained with the suppressed value of v_Δ even when the triplet fields are at the TeV scale. The Higgs sector is extended by introducing Z_2 -odd scalars (an $SU(2)_L$ doublet η and a real singlet s_2^0) in addition to a Z_2 -even complex singlet scalar s_1^0 whose vacuum expectation value violates the lepton number conservation by a unit. In our model, v_Δ is generated by the one-loop diagram to which Z_2 -odd particles contribute. The lightest Z_2 -odd scalar boson can be a candidate for the dark matter. We briefly discuss a characteristic signal of our model at the LHC.

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I. INTRODUCTION

Existence of dark matter (DM) has been established, and its thermal relic abundance has been determined by the WMAP experiment [1, 2]. If the essence of DM is an elementary particle, the weakly interacting massive particle (WIMP) would be a promising candidate. It is desired to have a viable candidate for the dark matter in models beyond the standard model (SM). The WIMP dark matter candidate can be accommodated economically by introducing only an inert scalar field [3–5], where we use “inert” for the Z_2 -odd property. The imposed Z_2 parity ensures the stability of the DM candidate. Phenomenology in such models have been studied in, e.g., Refs. [6–12].

On the other hand, it has been confirmed by neutrino oscillation measurements that neutrinos have nonzero but tiny masses as compared to the electroweak scale [13–17]. The different flavor structure of neutrinos from that of quarks and leptons may indicate that neutrino masses are of Majorana type. In order to explain tiny neutrino masses, many models have been proposed. The seesaw mechanism is the simplest way to explain tiny neutrino masses, in which right-handed neutrinos are introduced with large Majorana masses [18, 19]. Another simple model for generating neutrino masses is the Higgs Triplet Model (HTM) [19, 20]. However, these scenarios do not contain dark matter candidate in themselves.

In a class of models where tiny neutrino masses are generated by higher orders of perturbation, the DM candidate can be naturally contained [21–26]. In models in Refs. [21–25], the Yukawa couplings of neutrinos with the SM Higgs boson are forbidden at the tree level by imposing a Z_2 parity. The same Z_2 parity also guarantees the stability of the lightest Z_2 -odd particle in the model which can be the candidate of the DM as long as it is electrically neutral.

In this paper, we consider an extension of the HTM in which by introducing the Z_2 parity m_ν is generated at the one-loop level and the DM candidate appears. In the HTM, Majorana masses for neutrinos are generated via the Yukawa interaction $h_{\ell\ell'}\overline{L}_\ell^c i\sigma_2\Delta L_{\ell'}$ with a nonzero vacuum expectation value (VEV) of an $SU(2)_L$ triplet scalar field Δ with the hypercharge of $Y = 1$. The VEV of Δ is described by $v_\Delta \sim \sqrt{2}\mu v^2/(2M_\Delta^2)$, where v is the VEV of the Higgs doublet field Φ and M_Δ is the typical mass scale of the triplet field; the dimensionful parameter μ breaks lepton number conservation at the trilinear term $\mu\Phi^T i\sigma_2\Delta^\dagger\Phi$ which we refer to as the μ -term. As the simplest explanation for the smallness of neutrino masses, the

mass of the triplet field is assumed to be much larger than the electroweak scale. On the other hand a characteristic feature of the HTM is the fact that the structure of the neutrino mass matrix $(m_\nu)_{\ell\ell'}$ is given by that of the Yukawa matrix, $h_{\ell\ell'} \propto (m_\nu)_{\ell\ell'}$. The direct information on $(m_\nu)_{\ell\ell'}$ would be extracted from the decay $H^{\pm\pm} \rightarrow \ell^\pm \ell'^{\pm}$ [27] if H^{++} is light enough to be produced at collider experiments, where H^{++} is the doubly charged component of the triplet field Δ . At hadron colliders, the $H^{\pm\pm}$ can be produced via $q\bar{q} \rightarrow Z^*(\gamma^*) \rightarrow H^{++}H^{--}$ [28] and $q'\bar{q} \rightarrow W^{\pm*} \rightarrow H^{\pm\pm}H^\mp$ [29]. The $H^{\pm\pm}$ searches at the LHC put lower bound on its mass as $m_{H^{\pm\pm}} \gtrsim 300 \text{ GeV}$ [30, 31], assuming that the main decay mode is $H^{\pm\pm} \rightarrow \ell^\pm \ell'^\pm$. Phenomenological analyses for $H^{\pm\pm}$ in the HTM at the LHC have also been performed in Ref. [32]. Triplet scalars can contribute to lepton flavor violation (LFV) in decays of charged leptons, e.g., $\mu \rightarrow \bar{e}ee$ and $\tau \rightarrow \bar{\ell}\ell'\ell''$ at the tree level and $\ell \rightarrow \ell'\gamma$ at the one-loop level. Relation between these LFV decays and neutrino mass matrix constrained by oscillation data was discussed in Refs. [33, 34]. In order to explain the small v_Δ with such a detectable light H^{++} , the μ parameter has to be taken to be unnaturally much lower than the electroweak scale. Therefore, it would be interesting to extend the HTM in order to include a natural suppression mechanism of the μ parameter (therefore v_Δ) in addition to the DM candidate.

In our model, lepton number conservation is imposed to the Lagrangian in order to forbid the μ -term in the HTM at the tree level while the triplet Yukawa term $h_{\ell\ell'} \bar{L}_\ell^c i\sigma_2 \Delta L$ exists. The VEV of a Z_2 -even complex singlet scalar s_1^0 breaks the lepton number conservation by a unit. An $SU(2)_L$ doublet η and a real singlet s_2^0 are also introduced as Z_2 -odd scalars in order to accommodate the DM candidate. Then, the μ -term is generated at the one-loop level by the diagram in which the Z_2 -odd scalars are in the loop. By this mechanism, the smallness of $v_\Delta \ll v$ is realized, and the tiny neutrino masses are naturally explained without assuming the triplet fields to be heavy. The Yukawa sector is then the same as the one in the HTM, so that its predictions for the LFV processes are not changed. See Refs. [33, 35] for some discussions about two-loop realization of the μ -term¹.

This paper is organized as follows. In Sec. II, we give a quick review for the HTM to define notation. In Sec. III, the model for radiatively generating the μ parameter with the dark matter candidate is presented. Some phenomenological implications are discussed in

¹ The two-loop μ -term in Ref. [35] is given with softly-broken Z_4 symmetry, but the tree level μ -term would be also accepted as a soft breaking term. The two-loop μ -term in Ref. [33] is given with Z_3 symmetry which is broken by a VEV of a scalar S , but the tree level $\Phi^T i\sigma_2 \Delta^\dagger \Phi S^*$ seems allowed by the Z_3 .

Sec. IV, and the conclusion is given in Sec. V. The full expressions of the Higgs potential and mass formulae for scalar bosons in our model are given in Appendix.

II. HIGGS TRIPLET MODEL

In the HTM, an $SU(2)_L$ triplet of complex scalar fields with hypercharge $Y = 1$ is introduced to the SM. The triplet Δ can be expressed as

$$\Delta = \begin{pmatrix} \Delta^+/\sqrt{2} & \Delta^{++} \\ \Delta^0 & -\Delta^+/\sqrt{2} \end{pmatrix}, \quad (1)$$

where $\Delta^0 = (\Delta_r^0 + i\Delta_i^0)/\sqrt{2}$. The triplet has a new Yukawa interaction term with leptons as

$$\mathcal{L}_{\text{triplet-Yukawa}} = h_{\ell\ell'} \overline{L_\ell^c} i\sigma_2 \Delta L_{\ell'} + \text{h.c.}, \quad (2)$$

where $h_{\ell\ell'}$ ($\ell, \ell' = e, \mu, \tau$) are the new Yukawa coupling constants, $L_\ell [= (\nu_{\ell L}, \ell)^T]$ are lepton doublet fields, a superscript c means the charge conjugation, and σ_i ($i = 1-3$) denote the Pauli matrices. Lepton number ($L\#$) of Δ is assigned to be -2 as a convention such that the Yukawa term does not break the conservation. A vacuum expectation value $v_\Delta [= \sqrt{2} \langle \Delta^0 \rangle]$ breaks lepton number conservation by two units. The new Yukawa interaction then yields the Majorana neutrino mass term $(m_\nu)_{\ell\ell'} \overline{(\nu_{\ell L})^c} \nu_{\ell' L}/2$ where $(m_\nu)_{\ell\ell'} = \sqrt{2} v_\Delta h_{\ell\ell'}$.

The scalar potential in the HTM can be written as

$$\begin{aligned} V_{\text{HTM}} = & -m_\Phi^2 \Phi^\dagger \Phi + m_\Delta^2 \text{tr}(\Delta^\dagger \Delta) + \{\mu \Phi^T i\sigma_2 \Delta^\dagger \Phi + \text{h.c.}\} \\ & + \lambda_1 (\Phi^\dagger \Phi)^2 + \lambda_2 [\text{tr}(\Delta^\dagger \Delta)]^2 + \lambda_3 \text{tr}[(\Delta^\dagger \Delta)^2] \\ & + \lambda_4 (\Phi^\dagger \Phi) \text{tr}(\Delta^\dagger \Delta) + \lambda_5 \Phi^\dagger \Delta \Delta^\dagger \Phi, \end{aligned} \quad (3)$$

where $\Phi = (\phi^+, \phi^0)^T$ [$\phi^0 = (\phi_r^0 + i\phi_i^0)/\sqrt{2}$] is the Higgs doublet field in the SM. The μ parameter can be real by using rephasing of Δ . Because we take $m_\Delta^2 > 0$, there is no Nambu-Goldstone boson for spontaneous breaking of lepton number conservation. The small triplet VEV v_Δ is generated by an explicit breaking parameter μ of the lepton number conservation as

$$v_\Delta \simeq \frac{\sqrt{2} \mu v^2}{2m_\Delta^2 + (\lambda_4 + \lambda_5)v^2}, \quad (4)$$

where v ($\simeq 246$ GeV) is the doublet VEV defined by $v = \sqrt{2} \langle \phi^0 \rangle$.

In order to obtain small neutrino masses in the HTM, at least one of v^2/m_Δ^2 , $h_{\ell\ell'}$, μ/v should be tiny. A small μ is an attractive option because m_Δ can be small ($\lesssim 1$ TeV) so that triplet scalars can be produced at the LHC. Furthermore, large $h_{\ell\ell'}$ can be taken, which have direct information on the flavor structure of $(m_\nu)_{\ell\ell'}$. There is, however, no reason why the μ parameter is tiny in the HTM. In our model presented below, the μ parameter is naturally small because it arises at the one-loop level.

III. AN EXTENSION OF THE HIGGS TRIPLET MODEL

Since we try to generate the μ -term in the HTM radiatively, the term must be forbidden at the tree level. The simplest way would be to impose lepton number conservation to the Lagrangian. The conservation is assumed to be broken by the VEV of a new scalar field s_1^0 which is singlet under the SM gauge symmetry. Notice that $s_1^0 [= (s_{1r}^0 + is_{1i}^0)/\sqrt{2}]$ is a complex ("charged") field with non-zero lepton number although it is electrically neutral. One might think that the VEV of s_1^0 could be generated by using soft breaking terms of $L\#$. However, the μ -term is also a soft breaking term. Therefore lepton number must be broken spontaneously in our scenario. One may worry about Nambu-Goldstone boson corresponds to the spontaneous breaking of the lepton number conservation (the so-called Majoron, J^0). However the Majoron which comes from gauge singlet field can evade experimental searches (constraints) because it interacts very weakly with matter fields [36]. It is also possible to make it absorbed by a gauge boson by introducing the $U(1)_{B-L}$ gauge symmetry to the model (See, e.g., Ref. [37]). In this paper we just accept the Majoron without assuming the $U(1)_{B-L}$ gauge symmetry for simplicity.

If s_1^0 has $L\# = -2$, we can have a dimension-4 operator $\lambda s_1^0 \Phi^T i\sigma_2 \Delta^\dagger \Phi$. This gives a trivial result $\mu = \lambda \langle s_1^0 \rangle$ at the tree level. Although the dim.-4 operator could be forbidden by some extra global symmetries with extra scalars to break them, we do not take such a possibility in this paper. We just assume the s_1^0 has $L\# = -1$. Then the lepton number conserving operator which results in the μ -term is of dimension-5 as

$$(s_1^0)^2 \Phi^T i\sigma_2 \Delta^\dagger \Phi. \quad (5)$$

We consider below how to obtain the dim.-5 operator at the loop level by using renormal-

	L	Φ	Δ	s_1^0	s_2^0	η
$SU(2)_L$	<u>2</u>	<u>2</u>	<u>3</u>	<u>1</u>	<u>1</u>	<u>2</u>
$U(1)_Y$	1/2	1/2	1	0	0	1/2
$L\#$	1	0	-2	-1	0	-1
Z_2	+	+	+	+	-	-

TABLE I: List of particle contents of our one-loop model.

izable interactions². We restrict ourselves to extend only the $SU(3)_c$ -singlet scalar sector in the HTM because it seems a kind of beauty that the HTM does not extend the fermion sector and colored sector in the SM. An unbroken Z_2 symmetry is introduced in order to obtain dark matter candidates, and new scalars which appear in the loop diagram for the μ -term are aligned to be Z_2 -odd particles. We emphasize that the unbroken A_2 symmetry is not for a single purpose to introduce dark matter candidates but utilized also for our radiative mechanism for the μ -term.

We present the minimal model where the dim.-5 operator in eq. (5) is generated by a one-loop diagram with dark matter candidates. Table I shows the particle contents. A real singlet scalar field s_2^0 and the second doublet scalar field $\eta [= (\eta^+, \eta^0)^T, \eta^0 = (\eta_r^0 + i\eta_i^0)/\sqrt{2}]$ are introduced to the HTM in addition to s_1^0 . Lepton numbers of s_2^0 and η are 0 and -1, respectively. Then $\eta^T i\sigma_2 \Delta^\dagger \eta$ conserves lepton number. In order to forbid the VEV of η , we introduce an unbroken Z_2 symmetry for which s_2^0 and η are odd. Other fields are even under the Z_2 .

The Yukawa interactions are the same as those in the HTM. The Higgs potential is given as

$$V = \frac{1}{2}m_{s_2^0}^2(s_2^0)^2 + \{\mu_\eta \eta^T i\sigma_2 \Delta^\dagger \eta + \text{h.c.}\} + \{\lambda_{s\Phi\eta} s_1^0 s_2^0 (\eta^\dagger \Phi) + \text{h.c.}\} + \dots \quad (6)$$

Here we show only relevant parts for radiative generation of the μ -term. See Appendix for the other terms. Vacuum expectation values v and $v_s [= \sqrt{2} \langle s_1^0 \rangle]$ are given by

$$\begin{pmatrix} v^2 \\ v_s^2 \end{pmatrix} = \frac{2}{4\lambda_{1\Phi}\lambda_{s1} - \lambda_{s\Phi1}^2} \begin{pmatrix} 2\lambda_{s1} & -\lambda_{s\Phi1} \\ -\lambda_{s\Phi1} & 2\lambda_{1\Phi} \end{pmatrix} \begin{pmatrix} m_\Phi^2 \\ m_{s1}^2 \end{pmatrix}. \quad (7)$$

² It will not be difficult to do the same consideration for cases of higher dimensional operators, e.g., dim.-6 one with s_1^0 of $L\# = -2/3$.

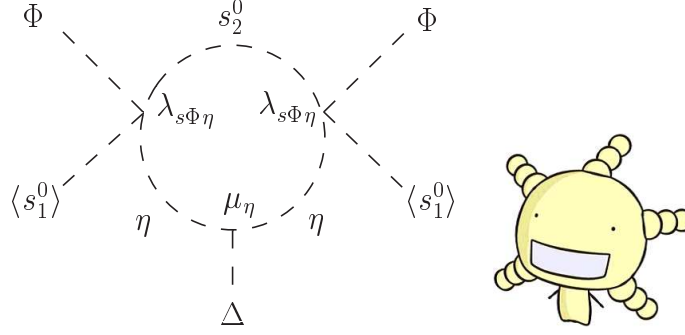


FIG. 1: One-loop diagram for the μ -term. We call it "A. oryzae diagram" [38].

The Z_2 -odd scalars in this model are two CP-even neutral ones (\mathcal{H}_1^0 and \mathcal{H}_2^0), a CP-odd neutral one ($\mathcal{A}^0 = \eta_i^0$), and a charged pair ($\mathcal{H}^\pm = \eta^\pm$). The CP-even scalars are defined as

$$\begin{pmatrix} \mathcal{H}_1^0 \\ \mathcal{H}_2^0 \end{pmatrix} = \begin{pmatrix} \cos \theta'_0 & -\sin \theta'_0 \\ \sin \theta'_0 & \cos \theta'_0 \end{pmatrix} \begin{pmatrix} \eta_r^0 \\ s_2^0 \end{pmatrix}, \quad \tan 2\theta'_0 = \frac{\sqrt{2} \lambda_{s\Phi\eta} v v_s}{(\mathcal{M}_0)_{ss}^2 - (\mathcal{M}_0)_{\eta\eta}^2}, \quad (8)$$

where $(\mathcal{M}_0)_{\eta\eta}^2 \equiv m_\eta^2 + (\lambda_{1\Phi\Phi} + \lambda_{1\Phi\eta}) v^2/2 + \lambda_{s\eta 1} v_s^2/2$ and $(\mathcal{M}_0)_{ss}^2 \equiv m_{s_2^0}^2 + \lambda_{s3} v_s^2 + \lambda_{s\Phi 2} v^2$.

Squared masses of these scalars are given by

$$m_{\mathcal{H}_1^0}^2 = \frac{1}{2} \left\{ (\mathcal{M}_0)_{\eta\eta}^2 + (\mathcal{M}_0)_{ss}^2 - \sqrt{\{(\mathcal{M}_0)_{\eta\eta}^2 - (\mathcal{M}_0)_{ss}^2\}^2 + 2 \lambda_{s\Phi\eta}^2 v^2 v_s^2} \right\}, \quad (9)$$

$$m_{\mathcal{H}_2^0}^2 = \frac{1}{2} \left\{ (\mathcal{M}_0)_{\eta\eta}^2 + (\mathcal{M}_0)_{ss}^2 + \sqrt{\{(\mathcal{M}_0)_{\eta\eta}^2 - (\mathcal{M}_0)_{ss}^2\}^2 + 2 \lambda_{s\Phi\eta}^2 v^2 v_s^2} \right\}, \quad (10)$$

$$m_{\mathcal{A}^0}^2 = (\mathcal{M}_0)_{\eta\eta}^2, \quad (11)$$

$$m_{\mathcal{H}^\pm}^2 = (\mathcal{M}_0)_{\eta\eta}^2 - \frac{1}{2} \lambda_{1\Phi\eta} v^2. \quad (12)$$

Notice that $m_{\mathcal{H}_1^0} \leq m_{\mathcal{A}^0} \leq m_{\mathcal{H}_2^0}$. We assume $m_{\mathcal{H}_1^0} < m_{\mathcal{H}^\pm}$ and then \mathcal{H}_1^0 becomes the dark matter candidate. Hereafter it is assumed that the mixing θ'_0 is small.

The μ -term is generated by the one-loop diagram. Figure 1 is the dominant one in the case of small θ'_0 . Then, the parameter μ is calculated as

$$\mu = \frac{\lambda_{s\Phi\eta}^2 \mu_\eta v_s^2}{64\pi^2 \{(\mathcal{M}_0)_{ss}^2 - (\mathcal{M}_0)_{\eta\eta}^2\}} \left\{ 1 - \frac{(\mathcal{M}_0)_{ss}^2}{(\mathcal{M}_0)_{ss}^2 - (\mathcal{M}_0)_{\eta\eta}^2} \ln \frac{(\mathcal{M}_0)_{ss}^2}{(\mathcal{M}_0)_{\eta\eta}^2} \right\}. \quad (13)$$

The one-loop induced μ parameter can be expected to be much smaller than μ_η . The suppression factor $|\mu/\mu_\eta|$ is estimated in Sec. IV A.

IV. PHENOMENOLOGY

A. Dark matter

If $(\mathcal{M}_0)_{\eta\eta} < (\mathcal{M}_0)_{ss}$, the dark matter candidate \mathcal{H}_1^0 is given by η_r^0 approximately because we assume small mixing. See, e.g., Ref. [8] for studies about the inert doublet scalar. Let us assume $m_{\mathcal{H}_1^0} \simeq 75 \text{ GeV}$ and $m_{\mathcal{A}^0} \gtrsim 125 \text{ GeV}$. As shown in Ref. [9], these values satisfy constraints from the LEP experiments [39, 40] and the WMAP experiment [2]. The mass splitting ($m_{\mathcal{A}^0} - m_{\mathcal{H}_1^0} \gtrsim 50 \text{ GeV}$) suppresses quasi-elastic scattering on nuclei ($\mathcal{H}_1^0 N \rightarrow \mathcal{A}^0 N$ mediated by the Z boson) enough to satisfy constraints from direct search experiments of the DM [41]. By using eqs. (9) and (11), we obtain

$$\frac{\lambda_{s\Phi\eta}^2 v_s^2}{(\mathcal{M}_0)_{ss}^2} \simeq \frac{2}{v^2} (m_{\mathcal{A}^0}^2 - m_{\mathcal{H}_1^0}^2) \gtrsim 0.3. \quad (14)$$

In order to be consistent with our assumption of small θ'_0 (e.g., $\simeq 0.1$), $(\mathcal{M}_0)_{ss} \gtrsim 3 \text{ TeV}$ is required. The value in eq. (14) results in

$$\frac{\mu}{\mu_\eta} \gtrsim 10^{-4}. \quad (15)$$

For the greater value of $m_{\mathcal{A}^0}$, the larger μ/μ_η is predicted. In particular, by taking $m_{\mathcal{A}^0}$ to be the TeV scale, we obtain $\mu/\mu_\eta \sim 10^{-2}$, which yields $v_\Delta \sim 1 \text{ GeV}$ for μ_η and m_Δ to be at the electroweak scale. Such a value for v_Δ is suggested in the recent study of radiative corrections to the electroweak parameters [42].

On the contrary, if we take $m_{\mathcal{A}^0} \simeq 83 \text{ GeV}$ which is allowed in a tiny region [9], values in eqs. (14) and (15) become 10 times smaller. We mention that the WMAP constraint might be changed by a characteristic annihilation process $\mathcal{H}_1^0 \mathcal{H}_1^0 \rightarrow \Delta_r^0 \rightarrow \bar{\nu} \bar{\nu}$ where $\mathcal{H}_1^0 \mathcal{H}_1^0 (\Delta_r^0)^*$ interaction is governed by μ_η (not by a tiny μ). This additional process could sift allowed value of $m_{\mathcal{H}_1^0}$ to lower one while $m_{\mathcal{A}^0} \gtrsim 100 \text{ GeV}$ due to the LEP constraint. Then, μ/μ_η might become larger than the value in eq. (15) because of larger $m_{\mathcal{A}^0} - m_{\mathcal{H}_1^0}$. This undesired effect would be easily avoided if $m_{\Delta_r^0}$ is away enough from $2m_{\mathcal{H}_1^0}$.

On the other hand, \mathcal{H}_1^0 comes dominantly from s_2^0 if $(\mathcal{M}_0)_{\eta\eta} > (\mathcal{M}_0)_{ss}$. See, e.g., Ref. [7] for studies about the real inert singlet scalar. Coupling $\sqrt{2} \lambda_{s\Phi 1} v$ of the $\mathcal{H}_1^0 \mathcal{H}_1^0 h^0$ interaction (h^0 is the SM Higgs boson) determines annihilation cross section of \mathcal{H}_1^0 and scattering cross section on nuclei. If we introduce the $U(1)_{B-L}$ gauge symmetry, the scattering of s_2^0 on nuclei can be mediated also by the gauge boson Z' . Notice that the parameter $\lambda_{s\Phi 1}$ (and

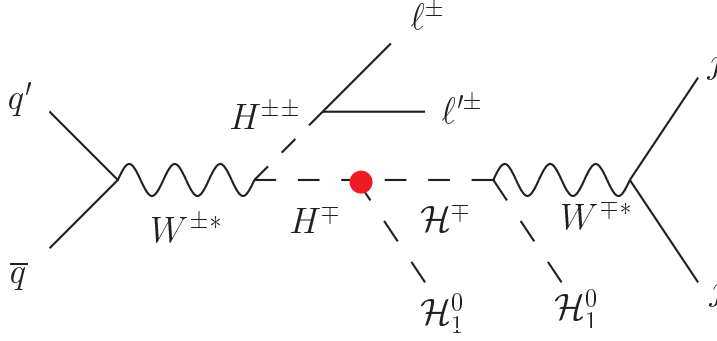


FIG. 2: The unique process in our model for $\mathcal{H}_1^0 \simeq \eta_r^0$. The bosonic decay of H^+ contains information of μ_η indicated by a red blob.

also the $U(1)_{B-L}$ gauge coupling constant) does not affect on μ parameter in eq. (13). Let us estimate the magnitude of μ/μ_η . In the usual HTM, $h_{\ell\ell'}$ is expected to be $\lesssim 10^{-2}$ for $m_{H^{\pm\pm}} \sim 100$ GeV in order to suppress LFV processes. Thus, we may accept $\lambda_{s\Phi\eta} \sim 1-10^{-2}$ as a value which is not too small. Assuming $(\mathcal{M}_0)_{ss} \ll (\mathcal{M}_0)_{\eta\eta} \sim v_s \sim 1$ TeV for example³, we have a suppression factor as

$$\left| \frac{\mu}{\mu_\eta} \right| \sim \frac{\lambda_{s\Phi\eta}^2 v_s^2}{64\pi^2 (\mathcal{M}_0)_{\eta\eta}^2} \sim 10^{-3}-10^{-7}. \quad (16)$$

Thus, even if the value of μ_η is in the TeV scale, we can obtain $\mu \sim 0.1$ MeV although we need further suppression with $h_{\ell\ell'} \lesssim 10^{-5}$ to have $m_\nu \lesssim 1$ eV. If we use $h_{\ell\ell'} \sim \lambda_{s\Phi\eta} \sim 10^{-3}$, we obtain $|\mu/\mu_\eta| h_{\ell\ell'} \sim 10^{-12}$ which can connect the TeV scale μ_η to the eV scale m_ν .

B. Collider

The characteristic feature of our model is that μ_η is much larger than μ . Let us consider possibility to probe the large μ_η in collider experiments.

A favorable process is shown in Fig. 2 for $\mathcal{H}_1^0 \simeq \eta_r^0$. For simplicity, we take $\lambda_5 = 0$ which results in $m_{H^{\pm\pm}} \simeq m_{H^\pm} \simeq m_{H^0, A^0}$. Recently, it was found in Ref. [42] that the electroweak precision test prefers $\lambda_5 > 0$ in the HTM where the electroweak sector is described by four input parameters. However, results in Ref. [42] might not be applied directly to our

³ If we introduce $U(1)_{B-L}$ gauge symmetry in order to eliminate the Majoron, v_s should be a little bit larger (e.g., ≥ 3 TeV) due to constraint on the mass of Z' .

model⁴ because the scalar sector is extended. Since $H^{\pm\pm} \rightarrow \ell^\pm \ell'^\pm$ is the most interesting decay in the HTM, we assume $2m_{\mathcal{H}^\pm} > m_{H^{\pm\pm}}$ in order to forbid $H^{\pm\pm} \rightarrow \mathcal{H}^\pm \mathcal{H}^\pm$. Even in this case, the DM \mathcal{H}_1^0 can be light enough ($m_{H^\pm} > m_{\mathcal{H}^\pm} + m_{\mathcal{H}_1^0}$) so that Z_2 -even charged scalar H^\pm ($\simeq \Delta^\pm$) can decay into $\mathcal{H}^\pm \mathcal{H}_1^0$ via μ_η -term which is indicated by a red blob in Fig. 2. The partial decay width of $H^\pm \rightarrow \mathcal{H}^\pm \mathcal{H}_1^0$ is determined by $(\mu_\eta/\mu)^2 v_\Delta^2/m_{H^\pm}$ while the width of $H^\pm \rightarrow \ell^\pm \nu$ is proportional to $m_{H^\pm} m_\nu^2/v_\Delta^2$. Taking $\mu_\eta/\mu \sim 10^4$, $v_\Delta \sim 10$ keV, $m_{H^\pm} \sim 100$ GeV, and $m_\nu \sim 0.1$ eV for example, we have $(\mu_\eta/\mu)^2 v_\Delta^2/m_{H^\pm} \sim 10^5$ eV and $m_{H^\pm} m_\nu^2/v_\Delta^2 \sim 10$ eV. Then, H^\pm dominantly decays into $\mathcal{H}^\pm \mathcal{H}_1^0$. Finally, \mathcal{H}^\pm decays into $(W^\pm)^* \mathcal{H}_1^0$. Therefore, from a production mechanism $pp \rightarrow (W^\pm)^* \rightarrow H^{\pm\pm} H^\mp$, we would have $\ell\ell jj \cancel{E}_T$ as a final state⁵ for which $\ell\ell$ has the invariant mass $m(\ell\ell)$ at $m_{H^{\pm\pm}}$ assuming that the value of $m_{H^{\pm\pm}}$ has been known already.

If $\mathcal{H}_1^0 \simeq s_2^0$, then H^\pm decays via μ_η -term into $\mathcal{H}^\pm \mathcal{A}^0$ or $\mathcal{H}^\pm \mathcal{H}_2^0$ followed by $\mathcal{H}_2^0 \rightarrow \mathcal{A}^0 J^0$ where a sizable $\lambda_{s\eta 1}$ is assumed⁶. Because of $\mathcal{A}^0 \rightarrow \mathcal{H}_1^0 J^0$ through $\lambda_{s\Phi\eta}$, we have again $\ell\ell jj \cancel{E}_T$ with $m(\ell\ell) = m_{H^{\pm\pm}}$ from $pp \rightarrow (W^\pm)^* \rightarrow H^{\pm\pm} H^\mp$.

In the usual HTM in contrast, the final state with such $\ell\ell$ is likely to include additional charged leptons ($\ell\ell\ell\ell$ from $H^{++}H^{--}$, $\ell\ell\ell\cancel{E}_T$ from $H^{\pm\pm}H^\mp$, etc.) if $H^{\pm\pm}$ decay dominantly into $\ell^\pm \ell'^\pm$. Therefore, our model would be supported if experiments observe final states which include jets and only two ℓ whose invariant mass gives $m(\ell\ell) = m_{H^{\pm\pm}}$. This potential signature might be disturbed by hadronic decays of τ because $H^{++}H^{--} \rightarrow \ell\ell\tau\tau$ can result in $\ell\ell jj \cancel{E}_T$ with $m(\ell\ell) = m_{H^{\pm\pm}}$. Realistic simulation is necessary to see the feasibility.

V. CONCLUSIONS AND DISCUSSION

We have presented the simple extension of the HTM by introducing a Z_2 -even neutral scalar s_1^0 of $L\# = -1$, a Z_2 -odd neutral real scalar s_2^0 of $L\# = 0$, and a Z_2 -odd doublet scalar field η of $L\# = -1$. The DM candidate \mathcal{H}_1^0 in our model is made from s_2^0 and η_r . The $\mu\Phi^T i\sigma_2 \Delta^\dagger \Phi$ interaction which is the origin of v_Δ (and neutrino masses) is induced at the one-loop level while the $\mu_\eta \eta^T i\sigma_2 \Delta^\dagger \eta$ interaction exists at the tree level. Because of the

⁴ Our model also has four parameters for the electroweak sector although v_Δ is generated at the 1-loop level.

⁵ Each of two \mathcal{H}_1^0 in Fig. 2 can be replaced with \mathcal{A}^0 which decays into $Z^* \mathcal{H}_1^0$ for $\mathcal{H}_1^0 \simeq \eta_r^0$.

⁶ If $\lambda_{s\eta 1}$ is small, \mathcal{H}_2^0 ($\simeq \eta_r^0$) decays into $Z^* \mathcal{A}^0$.

loop suppression for μ parameter, the model gives small neutrino masses naturally without using very heavy particles.

For $\mathcal{H}_1^0 \simeq \eta_r^0$, the suppression factor $|\mu/\mu_\eta|$ is constrained by the DM relic abundance measured by the WMAP experiment. We have shown that $|\mu/\mu_\eta| \sim 10^{-4} - 10^{-5}$ is possible. On the other hand, for $\mathcal{H}_1^0 \simeq s_2^0$, the suppression factor is somewhat free from experimental constraints on the DM. In our estimate, $|\mu/\mu_\eta| \sim 10^{-3} - 10^{-7}$ can be obtained as an example with $\lambda_{s\Phi\eta} \sim 1 - 10^{-2}$.

The characteristic feature of the model is that μ_η is not small while μ can be small. A possible collider signature which depends on μ_η would be $\ell\ell jj\cancel{E}_T$ with the invariant mass $m(\ell\ell) = m_{H^{\pm\pm}}$ because more charged leptons are likely to exist in such final states in the usual HTM.

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appendix

The Higgs potential of our model is given by $V = V_2 + V_3 + V_4$ where

$$V_2 \equiv -m_{s_1}^2 |s_1^0|^2 + \frac{1}{2} m_{s_2}^2 (s_2^0)^2 - m_\Phi^2 \Phi^\dagger \Phi + m_\eta^2 \eta^\dagger \eta + m_\Delta^2 \text{tr}(\Delta^\dagger \Delta), \quad (17)$$

$$V_3 \equiv (\mu_\eta \eta^T i\sigma_2 \Delta^\dagger \eta) + \text{h.c.}, \quad (18)$$

$$\begin{aligned}
V_4 \equiv & \lambda_{1\Phi} (\Phi^\dagger \Phi)^2 + \lambda_{1\eta} (\eta^\dagger \eta)^2 + \lambda_{1\Phi\Phi} (\Phi^\dagger \Phi)(\eta^\dagger \eta) + \lambda_{1\Phi\eta} (\Phi^\dagger \eta)(\eta^\dagger \Phi) \\
& + \lambda_2 [\text{tr}(\Delta^\dagger \Delta)]^2 + \lambda_3 \text{tr}[(\Delta^\dagger \Delta)^2] \\
& + \lambda_{4\Phi} (\Phi^\dagger \Phi) \text{tr}(\Delta^\dagger \Delta) + \lambda_{4\eta} (\eta^\dagger \eta) \text{tr}(\Delta^\dagger \Delta) \\
& + \lambda_{5\Phi} (\Phi^\dagger \Delta \Delta^\dagger \Phi) + \lambda_{5\eta} (\eta^\dagger \Delta \Delta^\dagger \eta) \\
& + \lambda_{s1} |s_1^0|^4 + \lambda_{s2} (s_2^0)^4 + \lambda_{s3} |s_1^0|^2 (s_2^0)^2 \\
& + \lambda_{s\Phi 1} |s_1^0|^2 (\Phi^\dagger \Phi) + \lambda_{s\Phi 2} (s_2^0)^2 (\Phi^\dagger \Phi) \\
& + \lambda_{s\eta 1} |s_1^0|^2 (\eta^\dagger \eta) + \lambda_{s\eta 2} (s_2^0)^2 (\eta^\dagger \eta) + \{ \lambda_{s\Phi\eta} s_1^0 s_2^0 (\eta^\dagger \Phi) + \text{h.c.} \} \\
& + \lambda_{s\Delta 1} |s_1^0|^2 \text{tr}(\Delta^\dagger \Delta) + \lambda_{s\Delta 2} (s_2^0)^2 \text{tr}(\Delta^\dagger \Delta).
\end{aligned} \tag{19}$$

All coupling constants are real because the phases of μ_η and $\lambda_{s\Phi\eta}$ can be absorbed by Δ and s_1^0 , respectively.

Mass eigenstates of two Z_2 -even CP-even neutral scalars which are composed of s_{1r}^0 and ϕ_r^0 are obtained as

$$\begin{pmatrix} h^0 \\ H^0 \end{pmatrix} = \begin{pmatrix} \cos \theta_0 & -\sin \theta_0 \\ \sin \theta_0 & \cos \theta_0 \end{pmatrix} \begin{pmatrix} \phi_r^0 \\ s_{1r}^0 \end{pmatrix}, \quad \tan 2\theta_0 = \frac{\lambda_{s\Phi 1} v v_s}{\lambda_{s1} v_s^2 - \lambda_{1\Phi} v^2}. \tag{20}$$

Their masses eigenvalues are given by

$$m_{h^0}^2 \simeq \lambda_{1\Phi} v^2 + \lambda_{s1} v_s^2 - \sqrt{(\lambda_{1\Phi} v^2 - \lambda_{s1} v_s^2)^2 + \lambda_{s\Phi 1}^2 v^2 v_s^2}, \tag{21}$$

$$m_{H^0}^2 \simeq \lambda_{1\Phi} v^2 + \lambda_{s1} v_s^2 + \sqrt{(\lambda_{1\Phi} v^2 - \lambda_{s1} v_s^2)^2 + \lambda_{s\Phi 1}^2 v^2 v_s^2}, \tag{22}$$

where small contributions from v_Δ are neglected. Two Z_2 -even CP-odd neutral bosons (ϕ_i^0 and s_{1i}^0) are Nambu-Goldstone bosons; ϕ_i^0 is absorbed by the Z boson, and s_{1i}^0 is the Majoron (or absorbed by the Z' boson).

Masses of bosons made dominantly from Δ are given by

$$m_{H_T^0}^2 \simeq m_{A_T^0}^2 \simeq m_{H^\pm}^2 + \frac{1}{4} \lambda_{5\Phi} v^2, \tag{23}$$

$$m_{H^\pm}^2 \simeq m_\Delta^2 + \frac{1}{4} (2\lambda_{4\Phi} + \lambda_{5\Phi}) v^2 + \frac{1}{2} \lambda_{s\Delta 1} v_s^2, \tag{24}$$

$$m_{H^{\pm\pm}}^2 \simeq m_{H^\pm}^2 - \frac{1}{4} \lambda_{5\Phi} v^2. \tag{25}$$

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